Loss of Andreev Backscattering in Superconducting Quantum Point Contacts

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We study effects of magnetic field on the energy spectrum in a superconducting quantum point contact. The supercurrent induced by the magnetic field leads to intermode transitions between the electron waves that pass and do not pass through the constriction. The latter experience normal reflections which couple the states with opposite momenta inside the quantum channel and create a minigap in the energy spectrum that depends on the magnetic field.

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The transport in superconducting–normal-metal (SN) mesoscopic hybrid structures is controlled by normal and Andreev reflections. The fundamental property of the Andreev reflection is its almost exact backscattering: trajectories of an incoming particle and reflected hole coincide within an angle of a maximum $\sim (k_F\xi)^{-1}$ where ξ is the superconducting coherence length and k_F is the Fermi wave vector. Backscattering holds also in the presence of magnetic field [1] since the electron-hole conversion length $\sim \xi$ is much smaller than the Larmor radius r_L, i.e., $\xi/r_L \sim (B/H_{c2})(k_F\xi)^{-1}$. This has been checked, for example, by the electron focusing technique [2]. Deviation from backscattering come from interference of electron and hole waves with an inhomogeneous order parameter phase. Effects of a supercurrent-induced transverse force on backscattering have been considered in Refs. [3, 4]. In Refs. [5] the phase difference introduced between the partial waves resulted in a decrease in the reflected hole intensity leading to an interplay between the normal and Andreev processes.

Though deviation from exact backscattering is small, it can still be noticeable if it is comparable with the size of the setup. A convenient device that satisfies this requirement is a ballistic superconducting quantum point contact. In the present paper we study how the violation of the fundamental semiclassical backscattering property of the Andreev reflection affects the subgap energy spectrum in such contact. Under the condition of exact backscattering, each electron that passes through the constriction is reflected as a hole that returns along the same trajectory and thus has also to pass through the constriction. The energy spectrum of subgap states is then $E = \pm\Delta_0 \cos\chi$ where Δ_0 is the superconducting gap and 2χ is the phase difference between the electrodes [6, 7]. The two energy branches which correspond to quasiparticles propagating through the constriction in opposite directions cross the Fermi level $E = 0$ at $2\chi = \pi$. We show that the loss of exact backscattering leads to a dramatic change in the spectrum such that a minigap appears near $2\chi = \pi$. The deviation from backscattering can be produced, for example, by an exchange, during the Andreev process, of a Cooper pair momentum induced by an applied magnetic field. This momentum mixes the channel modes with the modes that do not penetrate inside but are normally reflected from the channel end. The normal reflections couple the waves propagating through the constriction in the opposite directions and lead to formation of a minigap in the energy spectrum similar to that for contacts with normal scatterers [8, 9, 10, 11]. Varying the magnetic field one can tune the degree of normal reflection and manipulate the minigap thus controlling the transport properties of the contact.

Model.– The loss of exact backscattering at the Andreev reflections in a quantum point contact can be more clearly illustrated for an auxiliary structure shown in Fig. 1 (a): A single mode channel with a radius $a \sim k_F^{-1}$ is open into a normal semi-spherical region with a radius R much larger than the superconducting coherence length ξ. The normal region is surrounded by a superconductor which carries a supercurrent with a momentum $\hbar\mathbf{k}_s$ perpendicular to the channel axis. For $R \gg \xi$ one can describe quasiparticle propagation using a trajectory representation. Due to the transfer of $\hbar\mathbf{k}_s$ the Andreev reflected trajectory deviates from its initial direction [3] such that it can miss the constriction and experience normal reflections from the insulating barrier. The trajectory returns to the constriction after several reflections, which results in coupling of states propagating through the constriction in the opposite directions. The transfer of $\mathbf{k}_s$ causes a trajectory deflection by an angle k_s/k_F, thus the shift of the trajectory over a distance R would be k_sR/k_F. The probability of normal reflection thus depends on the ratio of the trajectory shift to the transverse channel dimension a. For a single-mode quantum channel $a \sim k_F^{-1}$, this ratio is $k_sR/k_Fa \sim k_sR$.

In a superconducting point contact, the quasiparticle wave functions for subgap states decay over distances of the order of $R \sim \xi$, which gives an estimate $k_sR \sim k_s\xi \lesssim 1$. However, the trajectory shift $k_s\xi/k_F$ is less than the wave length so that the trajectory pic-

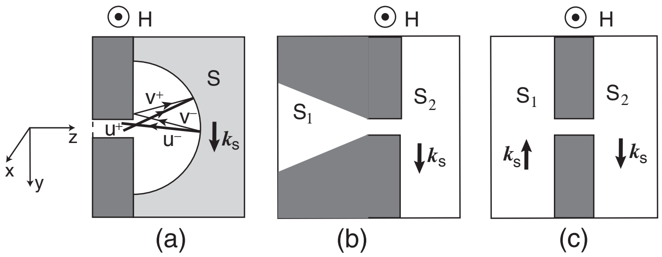

FIG. 1: (a) A single-mode channel is open to a normal region (white semicircle) in a contact with a superconductor (grey region). Andreev reflected trajectories deviate from initial direction due to the transverse pair momentum $\hbar\mathbf{k}_s$ and experience normal reflections from the insulator surface (black), which couple right-moving u^+ and left-moving u^- states. (b) Asymmetric and (c) symmetric point contacts.

ture is not applicable. Moreover, a single-mode contact with a radius $a \sim k_F^{-1}$ radiates an electronic wave which is determined by diffraction. To match the particle $u_0^\pm$ and hole $v_0^\pm$ waves propagating along the channel with the quasiparticle waves in the corresponding superconductors we assume, for simplicity, that the right superconductor occupies a halfspace $z > 0$ and has a specularly reflecting surface. We introduce the spherical coordinates $x = r\sin\theta\cos\phi, y = r\sin\theta\sin\phi, z = r\cos\theta$ and expand the wave functions in spherical functions with odd angular momenta. This ensures vanishing of the microscopic wave function at the superconducting/insulator boundary $\theta = \pi/2$. Far from the origin ($r \gg a$) we have

$$\begin{pmatrix} u^\pm \\ v^\pm \end{pmatrix} = \frac{e^{\pm ik_F r}}{r} \sum_{l=2n+1,m} P_{l,m}(\theta,\phi) \begin{pmatrix} U^\pm_{l,m}(r) \\ V^\pm_{l,m}(r) \end{pmatrix} . \quad (1)$$

Different angular harmonics are orthogonal to each other within $0 < \theta < \pi/2$. For a waveguide $a \lesssim k_F^{-1}$, the radiated/incident diffraction field is $\exp(\pm ik_F r)\cos\theta/r$. We now assume that it is this only mode proportional to $P_{1,0} = \cos\theta$ that ideally transforms into the channel mode u_0, v_0 without reflections, while all other modes $l \neq 1$ are normally reflected from the waveguide end without transmission into the channel.

Scattering Matrix.– For more quantitative description we use the Bogolubov - de Gennes (BdG) equations:

$$\left[\frac{1}{2m}\left(-i\hbar\boldsymbol{\nabla} - \frac{e}{c}\mathbf{A}\right)^2 - E_F\right] u + \Delta v \; = \; Eu \; , \quad (2)$$

$$\left[\frac{1}{2m}\left(-i\hbar\boldsymbol{\nabla} + \frac{e}{c}\mathbf{A}\right)^2 - E_F\right] v - \Delta^* u \; = \; -Ev \; , \quad (3)$$

where $\mathbf{A}$ is the vector potential of the magnetic field $\mathbf{B} = B(z)\hat{\mathbf{x}}$, and Δ is the gap function. We assume a step-like form $\Delta = \Delta_0 e^{i(\chi+\mathbf{k}_s\mathbf{r})}$, where $\chi_{R,L} = \pm\chi$ is the zero-field constant phase in the right (left) superconductor with 2χ phase difference between the superconductors; $\mathbf{k}_s$ is a constant wave vector. It enters the superconducting velocity $m\mathbf{v}_s = \hbar\mathbf{k}_s - (2e/c)\mathbf{A}$ that determines the difference in the eikonals of particle and hole wave functions, $(m/\hbar)\int \mathbf{v}_s \cdot d\mathbf{r}$. The magnetic field is screened in the superconductor, $\mathbf{A} = \lambda_L B_0 \hat{\mathbf{y}}(\exp(-z/\lambda_L) - 1)$, at the London length λ_L. From the screening condition $\mathbf{v}_s = 0$ at large z we find $\mathbf{k}_{sR} = (2\pi/\phi_0)\lambda_L B_0 \hat{\mathbf{y}}$, where ϕ_0 is the flux quantum. Assuming for simplicity $\lambda_L \gg \xi$ we neglect $\mathbf{A}$ in the region $r \sim \xi$ where the wave functions are localized. The parameter $k_s\xi$ that determines the probability of normal reflections of the channel modes is $k_s\xi \sim B_0/H_{cm}$ where $H_{cm} \sim \phi_0/(\lambda_L\xi)$ is the thermodynamic critical field of the superconductor. The gap Δ_0 is suppressed near the superconductor surface. However, this does not change the backscattering properties of Andreev reflection; we ignore it in what follows.

Inside the single-mode channel there are two particle and two hole waves $\propto e^{\pm ik_z z}$ with amplitudes $u_0^\pm$ and $v_0^\pm$, respectively, corresponding to the momentum projections $\pm\hbar k_z$ on the z axis. A particle u_0^+ and a hole v_0^- propagate in the $+z$ direction while a particle u_0^- and a hole v_0^+ propagate in $-z$ direction. Using the scheme employed in Ref. [10], we introduce the scattering matrices $\hat{S}_R(\epsilon,\chi,\mathbf{k}_{sR})$ and $\hat{S}_L(\epsilon,-\chi,\mathbf{k}_{sL})$ that relate the incident and outgoing wave amplitudes respectively at the right, $z = d/2$, and left, $z = -d/2$ ends of the channel:

$$\begin{pmatrix} u_0^- \\ v_0^+ \end{pmatrix}_R = \hat{S}_R \begin{pmatrix} u_0^+ \\ v_0^- \end{pmatrix}_R , \quad \begin{pmatrix} u_0^+ \\ v_0^- \end{pmatrix}_L = \hat{S}_L \begin{pmatrix} u_0^- \\ v_0^+ \end{pmatrix}_L . \quad (4)$$

Here $d \ll \xi$ is the channel length, $\mathbf{k}_{sL}$ and $\mathbf{k}_{sR}$ are the superflow momenta in the left and right superconductors, respectively. The wave functions at the both ends of the channel have different phase factors:

$$\begin{pmatrix} u_0^\pm \\ v_0^\pm \end{pmatrix}_R = e^{\pm ik_z d} \begin{pmatrix} u_0^\pm \\ v_0^\pm \end{pmatrix}_L . \quad (5)$$

The solvability condition of Eqs. (4) and (5) yields

$$\det\left(1 - e^{i\hat{\sigma}_z k_z d}\hat{S}_R e^{i\hat{\sigma}_z k_z d}\hat{S}_L\right) = 0 \; . \quad (6)$$

The matrix $\hat{S}$ is unitary: $\hat{S}\hat{S}^\dagger = 1$. Indeed, the BdG equations (2), (3) conserve the quasiparticle flow

$$\mathrm{div}\left[u^*\left(-i\hbar\boldsymbol{\nabla} - \frac{e}{c}\mathbf{A}\right)u + u\left(i\hbar\boldsymbol{\nabla} - \frac{e}{c}\mathbf{A}\right)u^* \right.$$
$$\left. -v^*\left(-i\hbar\boldsymbol{\nabla} + \frac{e}{c}\mathbf{A}\right)v - v\left(i\hbar\boldsymbol{\nabla} + \frac{e}{c}\mathbf{A}\right)v^*\right] = 0 \; .$$

Since this flow vanishes for $|E| < \Delta_0$ deep in the superconductor it should be zero also in the channel, whence $|u_0^+|^2 + |v_0^-|^2 = |u_0^-|^2 + |v_0^+|^2$ which results in $\hat{S}\hat{S}^\dagger = 1$.

We now calculate the matrix S explicitly with the account of reflections from the channel end. Wave functions decaying into the right superconductor obey the relations

$$v_R^+ \; = \; e^{-\frac{i}{2}(\chi+\mathbf{k}_{sR}\mathbf{r})}\check{a}_\epsilon^+(\mathbf{k}_{sR})e^{-\frac{i}{2}(\chi+\mathbf{k}_{sR}\mathbf{r})}u_R^+ \; , \quad (7)$$

$$u_R^- \; = \; e^{\frac{i}{2}(\chi+\mathbf{k}_{sR}\mathbf{r})}\check{a}_\epsilon^+(\mathbf{k}_{sR})e^{\frac{i}{2}(\chi+\mathbf{k}_{sR}\mathbf{r})}v_R^- \; , \quad (8)$$

that couple the electron and hole amplitudes near the channel end $|\mathbf{r}| \ll \xi$. Here

$$\check{a}^{\pm}_{\epsilon}(\mathbf{k}_s) = \epsilon + i\xi \mathbf{k}_s \boldsymbol{\nabla}/2k_F \mp i\sqrt{1-(\epsilon + i\xi \mathbf{k}_s \boldsymbol{\nabla}/2k_F)^2} \ ,$$

$\xi = \hbar v_F/\Delta_0$, and $\epsilon = E/\Delta_0$. For the left superconductor, similar expressions hold with $\chi \to -\chi$ and $\check{a}^+ \to \check{a}^-$.

Let us consider the matrix $\hat{S}$ at the right end of the channel. We place the origin of the coordinate system at the right channel end and consider the case $\chi = 0$ since, generally, $\hat{S}(\chi) = e^{i\chi\sigma_z/2}\hat{S}(\chi=0)e^{-i\chi\sigma_z/2}$. According to Eq. (1), the wave function amplitudes in the right superconductor at distances $a \ll r \ll \xi$ from the origin have the form (we omit the index R)

$$U^+ = u_0^+ P_{1,0} + \Psi_u^+ \ \ , \ \ V^+ = v_0^+ P_{1,0} + \Psi_v^+ \ ,$$
$$U^- = u_0^- P_{1,0} + \Psi_u \ \ , \ \ V^- = v_0^- P_{1,0} + \Psi_v \ .$$

The amplitudes Ψ_u and Ψ_v stand for the sum of components with $l \neq 1$ in Eq. (1) and describe the modes which experience normal reflections at the channel end. Introducing the operator $\check{R}_\epsilon$ that describes this reflection we can write the relation between the amplitudes $\Psi_u^+ = \check{R}_\epsilon \Psi_u$, $\Psi_v^+ = \check{R}_{-\epsilon}\Psi_v$. The functions Ψ_u, Ψ_v are orthogonal to the mode $P_{1,0}$: $\langle P_{1,0}|\Psi_{u,v}\rangle = 0$. The angular brackets denote the angular average within $0 < \theta < \pi/2$.

The modes Ψ_u and Ψ_v together with the traversing mode u_0, v_0 experience Andreev reflections while only u_0 and v_0 contribute to the flow through the channel. The unitarity $\hat{S}\hat{S}^\dagger = 1$ implies that the quasiparticles which are scattered normally off the superconductor surface and the channel end will eventually return into the constriction either as particles or as holes after certain number of Andreev reflections at the superconducting side. The Andreev relations (7,8) should be applied to the total functions U and V at a hemisphere with the center at the origin and the radius much smaller than ξ where $\mathbf{k}_s\mathbf{r} \sim k_s r \sin\theta \sin\phi \ll 1$. Taking the derivatives only of the rapidly varying exponents in Eq. (1) we obtain

$$v_0^+ P_{1,0} + \check{R}_{-\epsilon}\Psi_v \ = \ e^{-i\varphi_\epsilon}(u_0^+ P_{1,0} + \check{R}_\epsilon \Psi_u) \ , \quad (9)$$
$$u_0^- P_{1,0} + \Psi_u \ = \ -e^{i\varphi_{-\epsilon}}(v_0^- P_{1,0} + \Psi_v) \ , \quad (10)$$

where

$$e^{i\varphi_\epsilon} = \epsilon - \xi k_{sr}/2 + i\sqrt{1-(\epsilon - \xi k_{sr}/2)^2} \ , \quad (11)$$

and $k_{sr} = k_s \sin\theta\sin\phi$. Since the normal reflection at the channel end is associated with the momentum transfer of the order of $\hbar k_F$ one can neglect the dependence of $\check{R}$ on energy on the scale Δ_0 and assume $\check{R} = -e^{i\varphi_r}$ where φ_r is a constant phase shift. We then solve Eqs. (9), (10) for the functions Ψ_u, Ψ_v and use the orthogonality $\langle P_{1,0}|\Psi_{u,v}\rangle = 0$. This yields two equations that couple the amplitudes u_0^+, v_0^- to u_0^-, v_0^+ through the matrix

$$\hat{S} = \frac{1}{1-c_\epsilon^2}\left(e^{-i\varphi_r\hat{\sigma}_z}(|c_\epsilon|^2 - 1) - (c_\epsilon - c_\epsilon^*)e^{i\chi\hat{\sigma}_z}\hat{\sigma}_x\right) \quad (12)$$

where

$$c_\epsilon = \frac{\langle P_{1,0}\left(1+e^{i(\varphi_{-\epsilon}-\varphi_\epsilon)}\right)^{-1} P_{1,0}\rangle}{\langle P_{1,0}\left(e^{i\varphi_\epsilon}+e^{i\varphi_{-\epsilon}}\right)^{-1} P_{1,0}\rangle} \ . \quad (13)$$

One sees from Eq. (11) that $e^{i\varphi_\epsilon}(k_s) = -e^{-i\varphi_{-\epsilon}}(-k_s)$ whence $c_\epsilon^*(k_s) = -c_{-\epsilon}(-k_s)$. Moreover, $c_\epsilon(k_s)$ is an even function of k_s because a change in the sign of k_s can be compensated by the shift $\phi \to \pi + \phi$ in the integral over the angles. Thus, $c_\epsilon^*(k_s) = -c_{-\epsilon}(k_s)$. For small $k_s\xi$,

$$c_\epsilon \simeq e^{i\eta} - \frac{i(k_s\xi)^2\langle P_{1,0}^2 \sin^2\theta \sin^2\phi\rangle e^{2i\eta}}{8\langle P_{1,0}^2\rangle \sin^3\eta} \ . \quad (14)$$

Without k_s one has $c_\epsilon = e^{i\eta}$ where $e^{i\eta} = \epsilon + i\sqrt{1-\epsilon^2}$. As a result, the diagonal components of $\hat{S}$ vanish, thus the $+p_z$ and $-p_z$ states are decoupled.

In the diffraction picture, the transitions that couple the penetrating and non-penetrating modes are caused by the angle-dependent Doppler shift of energy proportional to k_{sr} in Eq. (11). As a result, the wave fronts of reflected holes are distorted as compared to those of incident particles. The interference of these waves near the channel end results in the suppression of the amplitude of the Andreev reflected wave entering the channel.

Results.– Consider first the zero-bias conductance of a normal-metal – quantum-channel–superconductor junction [12] $G_s = (e^2/\pi\hbar)(1-|S_{11}|^2 + |S_{12}|^2)$ where $|S_{11}|^2$ and $|S_{12}|^2$ are probabilities of normal and Andreev reflection, respectively. We get for small $k_s\xi$

$$G_s = \frac{e^2}{\pi\hbar}\left[2 - \frac{2(|c_\epsilon|^2-1)^2}{|c_\epsilon^2 - 1|^2}\right]_{\epsilon=0} \simeq \frac{e^2}{\pi\hbar}\left[2 - \frac{1}{2}\left(\frac{B_0}{H_c}\right)^4\right] \ .$$

Here we introduce a field $H_c \sim H_{cm}$ through

$$\frac{B_0^2}{H_c^2} = \frac{(k_s\xi)^2\langle P_{1,0}^2\sin^2\theta\sin^2\phi\rangle}{4\langle P_{1,0}^2\rangle} = \frac{(k_s\xi)^2}{20} \ . \quad (15)$$

Consider now an asymmetric structure that consists of a superconducting tip with a curvature radius smaller than λ_L in a contact with a bulk superconductor, see Fig. 1 (b). In this case $\mathbf{k}_{sL} = 0$ while $\mathbf{k}_{sR} = \mathbf{k}_s \neq 0$. On the right end of the channel the matrix $\hat{S}_R = \hat{S}(\epsilon, \chi, \mathbf{k}_s)$ is determined by Eq. (12). On the left end the matrix $\hat{S}_L = \hat{S}(\epsilon, -\chi, 0)$ assumes an Andreev form $\hat{S}_L = e^{-i\eta}e^{-i\chi\sigma_z}\hat{\sigma}_x$. The phase shift $k_z d - \varphi_r$ drops out and Eq. (6) yields

$$\left(1-c_\epsilon^2\right)e^{i\eta} - \left(1-c_\epsilon^{*2}\right)e^{-i\eta} = 2\left(c_\epsilon^* - c_\epsilon\right)\cos(2\chi) \ . \quad (16)$$

For $k_s = 0$ and $c_\epsilon = e^{i\eta}$ we obtain a standard gapless expression $\epsilon = \pm\cos\chi$. For a nonzero k_s, a gap opens in the energy spectrum. Indeed, consider Eq. (16) in the limit of small magnetic fields and energies. It becomes

$$\epsilon^2 = \cos^2\chi + \frac{1}{8}\left[(ic_\epsilon + 1)^2 + (ic_\epsilon^* - 1)^2\right]_{\epsilon=0} \ .$$

Within the leading terms in B/H_c we find

$$\epsilon^2 = \cos^2\chi + \epsilon_g^2 \ . \tag{17}$$

where the minigap in the spectrum is $\epsilon_g = \frac{1}{4}(B_0/H_c)^2$.

In the case of a symmetric contact shown in Fig. 1 (c) the solution of the screening problem yields $\mathbf{k}_{sL} = -\mathbf{k}_{sR} = -\mathbf{k}_s$. The spectral equation (6) with $\hat{S}_R = \hat{S}(\epsilon, \chi, \mathbf{k}_s)$ and $\hat{S}_L = \hat{S}(\epsilon, -\chi, -\mathbf{k}_s)$ reduces to

$$\frac{(c_\epsilon + c_\epsilon^*)^2}{4} = \cos^2\chi + \frac{(|c_\epsilon|^2-1)^2}{(ic_\epsilon - ic_\epsilon^*)^2}\sin^2(k_z d - \varphi_r) \ . \tag{18}$$

The spectrum has a gap in the presence of k_s. In the limit of low energies and small magnetic fields the right hand side of Eq. (18) can be treated as a perturbation. We put $(|c_\epsilon|^2-1)^2 \simeq (B_0/H_c)^4$ where H_c is defined in Eq. (15) while $(c_\epsilon - c_\epsilon^*)^2 \approx -4$. At the same time, $(c_\epsilon + c_\epsilon^*)^2/4 \simeq \epsilon^2$. Finally we get Eq. (17) where

$$\epsilon_g = \frac{1}{2}(B_0/H_c)^2 \left|\sin(k_z d - \varphi_s)\right| \ . \tag{19}$$

Discussion.– Since the wave vector $k_s \lesssim \xi^{-1}$ is much smaller than k_F it induces transitions only between the modes with close transverse quantum numbers. Thus, the predicted effect can be more easily seen for a contact that is transparent only for a few modes. On the contrary, for a multi-mode channel, the coupling to the reflected modes that mixes $\mathbf{p}$ and $-\mathbf{p}$ states has a small weight while the transitions occur mostly between the penetrating modes. For a large area SNS junction these transitions result in the subgap spectrum instability with formation of energy bands [1].

Equation (17) coincides with the spectrum in the presence of normal scatterers [10, 11]. Note that the gap in a symmetric contact Eq. (19) vanishes for certain phase difference $k_z d - \varphi_r = \pi n$. This is a result of resonant tunneling through a system of two barriers with equal reflection coefficients. The transmission probability $|T|^2$ through such system is unity at the resonance such that the gap $\epsilon_g = 1 - |T|^2$ disappears. A small asymmetry in the scattering conditions removes the resonant tunneling effects so that a gap will exist for any phase shift $k_z d - \varphi_r$. The extreme asymmetric case is illustrated by Eq. (17). Similar effects of resonant tunneling and minigap oscillations as functions of $k_z d$ can also take place for other mechanisms of normal reflection such as interface barriers, mismatch in the material parameters, or small normal scattering from the step-like gap potential [13].

The predicted spectrum can be tuned by varying the magnetic field. The minigap is not small and can reach values of the order of Δ_0 for $B_0 \sim H_{cm}$. It can be monitored by measuring the Josephson critical current that decreases with the minigap [11]. Moreover, the minigap affects the dynamic properties of the point contact. In particular, it is responsible for suppression of the time-averaged quasiparticle current for voltage biased contacts in the region $eV < \epsilon_g$ [14]. Varying the magnetic field one can thus observe a transition from ballistic to high-resistance behavior of the contact.

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